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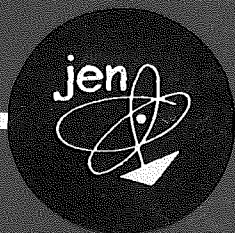
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**Self-similar motion of laser half-space
plasmas. II. Thermal wave
and intermediate regimes.**

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CLASIFICACION INIS Y DESCRIPTORES

A14

PLASMA WAVES

HOT PLASMA

LASER PRODUCED PLASMA

SHOCK WAVES

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NOTA

El presente trabajo corresponde al Informe #3 del equipo ii) Propagación de Ondas de Choque en Medios Ultradensos, del Subprograma de la JEN sobre Confinamiento Inercial del Proyecto Laser-Fisión-Fusión, descrito en el Informe JEN 351.

I. INTRODUCTION

In a previous paper ¹ (called I hereafter) the authors analyzed the self-similar motion generated in a fully ionized, unmagnetized plasma (initially cold and occupying the half-space $x>0$) by a laser pulse of irradiation $\phi \equiv \phi_0 t/\tau$ ($0 < t \leq \tau$), uniformly absorbed in the plane where the electron density equals the critical density ($n_e = n_c \equiv m_e \omega^2 / 4\pi e^2$, ω being the laser frequency). The equations governing the problem were found to depend on three dimensionless numbers: the ion charge number Z_i ; a parameter $\alpha \propto (n_0^2 \tau / \phi_0)^{2/3}$, where n_0 is the undisturbed electron density;^{2,3} and the ratio $n_c / n_0 \equiv \epsilon$ (assumed to be small). Detailed results were given in I for the regime $\alpha \gg \epsilon^{-4/3}$.

Here we study the regimes $1 < \alpha < \epsilon^{-4/3}$ and $\alpha \ll 1$. We notice that defining $\alpha_c \equiv \alpha \epsilon^{4/3} \propto (n_c^2 \tau / \phi_0)^{2/3}$, ($\alpha_c \ll \alpha$), the three main regimes just mentioned correspond to conditions $1 < \alpha_c < \alpha$, $\alpha_c \ll 1 < \alpha$, and $\alpha_c \ll \alpha \ll 1$ respectively.

Section II briefly reviews the equations established in I. Sections III and IV deal with the regimes $\alpha \ll 1$ and $1 < \alpha < \epsilon^{-4/3}$ respectively. In Sec. V we present a discussion of some results obtained and of the assumptions used in the analysis. Finally, Sec. VI carries out a qualitative discussion of how the plasma behavior changes when α sweeps from small to larger than $\epsilon^{-4/3}$ values through the intermediate regime $1 < \alpha < \epsilon^{-4/3}$ and the transition ranges $\alpha \sim 1$, $\alpha \sim \epsilon^{-4/3}$.

II. STATEMENT OF THE PROBLEM

Following I, we assume quasineutrality, so that $n_e \approx n_i Z_i \equiv n$, and $v_e \approx v_i \equiv v$ (since the current density must vanish in the undisturbed plasma); we retain, however, different temperatures for each species. We also assume (to start with) that the plasma is dominated by collisions, so that the electron heat conductivity and the ion-electron energy relaxation time may be written as.⁴

$$K_e = \bar{K}_e T_e^{5/2}, \quad t_{ei} = \bar{t}_{ei} T_e^{3/2} / n_e,$$

where \bar{K}_e and \bar{t}_{ei} are constant since we neglect the variations in all Coulomb logarithms; viscosities and ion conductivity are found numerically negligible.¹ As in I, we introduce self-similar variables

$$\xi = x / [w\tau(t/\tau)^{4/3}], \quad u(\xi) = v / [v_0(t/\tau)^{1/3}] \quad (1)$$

$$\bar{n}(\xi) = n/n_0, \quad \theta_j = T_j / [T_0(t/\tau)^{2/3}], \quad (j=e, i),$$

and arrive at the equations

$$\frac{d\bar{n}}{d\xi} = \frac{\bar{n}}{\xi-u} \frac{du}{d\xi}, \quad (2)$$

$$u - 4(\xi - u) \frac{du}{d\xi} = -\frac{\alpha}{\bar{n}} \frac{d}{d\xi} [\bar{n}(Z_i \theta_e + \theta_i)], \quad (3)$$

$$\bar{n} \left[\theta_e \left(1 + \frac{4}{3} \frac{du}{d\xi} \right) - 2(\xi - u) \frac{d\theta_e}{d\xi} \right] =$$

$$= \frac{d}{d\xi} \left(\chi \theta_e^{5/2} \frac{d\theta_e}{d\xi} \right) - 4.3 b(Z_i) \alpha \bar{n}^{-2} \frac{\theta_e - \theta_i}{\theta_e^{3/2}} + \delta(\xi - \xi_c), \quad (4)$$

$$\frac{\bar{n}}{Z_i} \left[\theta_i \left(1 + \frac{4}{3} \frac{du}{d\xi} \right) - 2(\xi - u) \frac{d\theta_i}{d\xi} \right] = 4.3 b(Z_i) \alpha \bar{n}^{-2} \frac{\theta_e - \theta_i}{\theta_e^{3/2}}, \quad (5)$$

$$\bar{n}(\xi_c) = n_c/n_o \equiv \epsilon \quad (6)$$

Electron inertia has been neglected in the momentum equation for the ion-electron fluid, from which the electric field has disappeared because of quasineutrality.⁴ Expressions for w , v_o , T_o , α and b are given in I.

The boundary conditions are: in the undisturbed plasma ($\xi=\infty$)

$$u = \theta_e = \theta_i = 0, \quad \bar{n} = 1; \quad (7)$$

at the vacuum-plasma boundary ($\xi=\xi_v$)

$$\bar{n} = 0, \quad u = \xi_v. \quad (8)$$

In addition, since the mean free path is proportional to θ_e^2/\bar{n} , θ_e should vanish at ξ_v for a collision-dominated plasma. We advance, however, that, for the regimes here considered, the plasma will be found to be collisionless near ξ_v . Thus, we first simply require that the heat flux vanish at ξ_v , and second, include a heat flux limiter χ in Eq. (4)^{5,6}

$$\chi = \left[1 + \left(\frac{8\pi m_e}{9k^3} \right)^{1/2} \frac{\bar{K}_e}{n_e} T_e \left| \frac{dT_e}{dx} \right| \right]^{-1} = \left[1 + \frac{0.059}{\alpha^{1/2} A_i^{1/2}} \frac{\theta_e}{\bar{n}} \left| \frac{d\theta_e}{d\xi} \right| \right]^{-1}, \quad (9)$$

where k is Boltzmann's constant and A_i the ion mass number.

III. ASYMPTOTIC SOLUTION FOR $\alpha \ll 1$

A. Thermal Wave Region

In Ref. 3, the motion resulting from deposition of energy at the plane $\xi=0$, within an unbounded, initially cold and

uniform, plasma, was considered; for α small, a thermal wave was found to move into the undisturbed plasma, convection and ion temperature being negligible. Following the analysis there we expand the dependent variables in Eqs. (2)-(5) in powers of α and find to lowest order

$$u \sim \bar{n}^{-1} \sim \theta_i \sim 0(\alpha), \quad \theta_e \sim 0(1),$$

the equation for θ_e reading

$$\theta_e^{-2} \xi d\theta_e/d\xi = d(\theta_e^{5/2} d\theta_e/d\xi)/d\xi, \quad (10)$$

where we assumed $\chi \approx 1$.

Since the front ξ_f of the thermal wave must lie at a finite distance,⁷ boundary conditions at $\xi = \infty$ are useless; however, vanishing of both electron heat flux and temperature at ξ_f are now required. In addition we have

$$\theta_e^{5/2} d\theta_e/d\xi = -1 \quad \text{at } \xi = 0^+ \quad (11)$$

since, as we shall see, the thickness of the expanding plasma in the negative- ξ half-space is $0(\alpha^{1/2})$ and therefore most of the energy is carried in the thermal wave; then, first, in a scale where $\xi = 0(1)$ the critical plane lies at the origin, and second, the leftward heat flux at that plane is negligible. The appropriate solution to Eq. (10) is given in Fig. 1, and differs from that in Ref. 3 because of condition (11).

To dominant order, Eqs. (2), (3), and (5) read

$$d\bar{n}/d\xi = \xi^{-1} du/d\xi \quad (12)$$

$$u - 4\xi du/d\xi = -\alpha Z_1 d\theta_e/d\xi \quad (13)$$

$$\theta_i - 2\xi d\theta_i/d\xi = 4.3 \alpha Z_i b \theta_e^{-1/2} \quad (14)$$

with boundary conditions

$$u = \bar{n} - 1 = \theta_i = 0 \quad \text{at} \quad \xi = \xi_f. \quad (15)$$

Solutions to Eqs. (12)-(15) are given in Fig. 1 for $\xi > 0$.

We notice that at $\xi=0$ both u and θ_i present a weak discontinuity (cusp) while $\bar{n}-1$ goes to infinity. Actually, when $\xi=0(\alpha^{1/2})$ the pressure gradient due to density variations, neglected in (13), must be retained. The analysis of this thin layer would be then similar to that carried out in Ref. 3: The cusp is found to move to $\xi=(1/2)[\theta_e(0)Z_i\alpha]^{1/2}$. To its right u is determined by matching to the solution in Fig. 1; its value at the cusp is $u(\xi=0)$ as given in that figure, while $\bar{n}-1$ reaches there a finite value of order $\alpha^{5/8}$. On the other hand, the behavior to the left of the cusp depends on the boundary conditions at ξ_v , and bears no resemblance to results in Ref. 3.

B. Isothermal Expansion Region

Boundary conditions (8) imply that u and ξ are of the same order, while $\bar{n}=0(1)$. It follows clearly from Eq. (4) that

$$\theta_e \approx \text{constant} = \theta_e(0).$$

Defining now

$$\hat{u} = u / [\alpha \theta_e(0) Z_i]^{1/2} \quad (16)$$

$$\hat{\xi} = \xi / [\alpha \theta_e(0) Z_i]^{1/2} \quad (17)$$

$$\hat{\theta}_i = \theta_i / \alpha b Z_i \quad (18)$$

Eqs. (2), (3), and (5) become

$$\frac{d\bar{n}}{d\hat{\xi}} = \frac{\bar{n}}{\hat{\xi}-\hat{u}} \frac{d\hat{u}}{d\hat{\xi}} \quad (19)$$

$$\hat{u}-4(\hat{\xi}-\hat{u}) \frac{d\hat{u}}{d\hat{\xi}} = \frac{-1}{\bar{n}} \frac{d\bar{n}}{d\hat{\xi}} \quad (20)$$

$$\hat{\theta}_i \left(1 + \frac{4}{3} \frac{d\hat{u}}{d\hat{\xi}}\right) - 2(\hat{\xi}-\hat{u}) \frac{d\hat{\theta}_i}{d\hat{\xi}} = \frac{4.3 \bar{n}}{[\theta_e(0)]^{1/2}} \quad (21)$$

The boundary conditions are

$$\bar{n}=0, \quad \hat{u}=\hat{\xi}_v \quad \text{at} \quad \hat{\xi}=\hat{\xi}_v \quad (22)$$

$$\bar{n}=1, \quad \hat{u}=0, \quad \hat{\theta}_i = \frac{\theta_i(\text{cusp})}{\alpha b z_i} \approx 4.27 \quad \text{at} \quad \hat{\xi}=\frac{1}{2}, \quad (23)$$

the value of θ_i at the cusp clearly being $\theta_i(\xi=0)$ as given in Fig. 1; on the other hand, the conditions on \hat{u} and \bar{n} at $\hat{\xi}=1/2$ follow from the fact that at the cusp $\hat{u}=0(\alpha^{1/2})$ and $\bar{n}-1=0(\alpha^{5/8})$.

We now notice that once \hat{u} and \bar{n} have been obtained $\hat{\theta}_i$ follows by direct integration. Moreover, a single equation for \hat{u} may be immediately obtained from (19) and (20):

$$\frac{d\hat{u}}{d\hat{\xi}} = \frac{-\hat{u}(\hat{\xi}-\hat{u})}{1-4(\hat{\xi}-\hat{u})^2} \quad (24)$$

$$\hat{u}=0 \quad \text{at} \quad \hat{\xi}=1/2. \quad (25)$$

It should be pointed out, however, that $\hat{\xi}=1/2, \hat{u}=0$ is a nodal point of (24), and therefore there is an infinite number of solutions to Eqs. (24)-(25). Now, for a particular solution, Eq. (19) may be integrated, starting at $\hat{\xi}=1/2, \bar{n}=1$: clearly the so-

lution looked for is the one for which $\bar{n}=0$ when $\hat{u}=\hat{\xi}$. We find finally that this solution meets the line $\hat{u}=\hat{\xi}$ at $\hat{\xi}=-\infty$; all other solutions to Eqs. (24)-(25) either cross that line at a finite $\hat{\xi}$, where $\bar{n}\neq 0$, or do not meet it at all.

Figure 2 shows then all three \hat{u} , \bar{n} , and $\hat{\theta}_i$ throughout the expansion layer; although, asymptotically, $\hat{\xi}_c \rightarrow -\infty$ as $\epsilon \rightarrow 0$, $\bar{n}(\hat{\xi})$ decreases so fast that, for any reasonable ϵ , $\hat{\xi}_c$ has a moderate value. Actually, as we shall see in Sec. V, the solution shown in Fig. 2 breaks down when \bar{n} becomes small because the plasma becomes collisionless.

IV. ASYMPTOTIC SOLUTION FOR $1 < \alpha < \epsilon^{-4/3}$

A. Isentropic Compression Region

It was shown in I that as long as α is large there exists a region of isentropic compression behind a shock bounding the undisturbed plasma; its analysis would be here entirely similar to that carried out in I. In fact, if $\bar{n}_f=4$, $u_f=3\xi_f/4$, and $\theta_f=3\xi_f^2/4\alpha(Z_i+1)$, are the values behind the shock, as obtained from the jump conditions (ξ_f being the unknown shock position), the solution given in figure 1 there in normalized variables

$$\begin{aligned} \eta &\equiv \xi/\xi_f, & v &\equiv \bar{n}/\bar{n}_f \\ y &\equiv u/u_f, & z_j &\equiv \theta_j/\theta_f \quad (z_i \approx z_e \equiv z), \end{aligned} \quad (26)$$

remains valid for the present regime. It was found in I that the isentropic solution ceases to be valid near a point $\bar{\eta}(\bar{\eta} \approx 0.82)$, where

$$y \approx \frac{4}{3} \bar{\eta} - \frac{2}{5} (\eta - \bar{\eta}), \quad z \approx B_1 (\eta - \bar{\eta})^{3/13}, \quad v \approx B_2 (\eta - \bar{\eta})^{-3/13}$$

($B_1 \approx 1.70$, $B_2 \approx 0.78$), the plasma becoming there highly dense and cold.

B. Expansion Region

We shall now make the ansatz that a) to analyze the entire region between η_v and $\bar{\eta}$ it suffices to scale the variables once, and that b) $v \leq 0(1)$. We then notice first, that

$$vz = 0(1) \quad (27)$$

since it must have the value $B_1 B_2$ at $\bar{\eta}$, and secondly that condition (8) implies

$$y = 0(\eta) \quad (28)$$

Next, integrating Eq. (4) across a thin layer centered at the critical plane, and using the variables defined in (26), we get

$$\eta = 0(\alpha^{-7/2} \xi_f^6 z^{7/2}) \quad (29)$$

In addition, the condition of vanishing momentum for the entire plasma leads to

$$v\eta = 0(1) \quad (30)$$

since in the compression region $\int_{\bar{\eta}}^1 v y d\eta = 0(1)$. Finally the integral energy conservation law, which follows from system (2)-(5), may be written as

$$\int_{\xi_v}^{\xi_f} 3 \bar{n} (Z_i \theta_e + \theta_i + 4u^2 / 3\alpha) d\xi = Z_i ;$$

this equation, using (26) and $v \leq 0(1)$, shows that

$$\xi_f^3 \eta = 0(\alpha) \quad . \quad (31)$$

Equations (27)-(31) lead to $\xi_f = 0(\alpha^{5/24})$ and suggest defining new variables

$$\begin{aligned} \hat{\eta} &= \xi_f^{3/5} \alpha^{-1/2} \eta, & \hat{y} &= \xi_f^{3/5} \alpha^{-1/2} y \\ \hat{v} &= \xi_f^{-6/5} \alpha v, & \hat{z}_j &= \xi_f^{6/5} \alpha^{-1} z_j \end{aligned} \quad (32)$$

Then, Eqs. (2)-(5) become

$$\frac{d\hat{v}}{d\hat{\eta}} = \frac{3\hat{v}}{4\hat{\eta}-3\hat{y}} \frac{d\hat{y}}{d\hat{\eta}} \quad (33)$$

$$\hat{v}\hat{y} - \hat{v}(4\hat{\eta}-3\hat{y}) \frac{d\hat{y}}{d\hat{\eta}} = - \frac{d}{d\hat{\eta}} \left(\hat{v} \frac{z_i \hat{z}_e + \hat{z}_e}{z_i + 1} \right) \quad (34)$$

$$\begin{aligned} \frac{\hat{v}}{4} \left[\hat{z}_e \left(1 + \frac{d\hat{y}}{d\hat{\eta}} \right) - \frac{4\hat{\eta}-3\hat{y}}{2} \frac{d\hat{z}_e}{d\hat{\eta}} \right] &= \frac{1}{16} \left[\frac{3}{4(Z_i+1)} \right]^{5/2} \frac{d}{d\hat{\eta}} \hat{z}_e^{5/2} \frac{d\hat{z}_e}{d\hat{\eta}} - \\ &- 4.3 b \left[\frac{4(Z_i+1)}{3} \right]^{3/2} \hat{v}^2 \frac{\hat{z}_e - \hat{z}_i}{\hat{z}_e^{3/2}} \end{aligned} \quad (35)$$

$$\frac{\hat{z}_i}{4} \left(1 + \frac{d\hat{y}}{d\hat{\eta}} \right) - \frac{4\hat{\eta}-3y}{8} \frac{d\hat{z}_i}{d\hat{\eta}} = 4.3 b \left[\frac{4(Z_i+1)}{3} \right]^{3/2} \hat{v} \frac{\hat{z}_e - \hat{z}_i}{\hat{z}_e^{3/2}} ; \quad (36)$$

explicit use of ξ_f in (32) has allowed to drop it from Eqs. (33)-(36).

This system must be solved subject to five boundary conditions: these are

$$\hat{y} = \hat{z}_e = \hat{z}_i = 0, \quad \hat{v}\hat{z}_e = B_1 B_2 \quad \text{at} \quad \hat{\eta} = \xi_f^{3/5} \alpha^{-1/2} \bar{\eta} \approx 0, \quad (37)$$

(which follow from matching to the isentropic compression solution), together with the vacuum-plasma boundary condition,

$$\bar{n}=0 \quad \text{when} \quad 4\hat{\eta}=3\hat{y} \quad . \quad (38)$$

There is no δ -function term in (35) because $\hat{v}_c=0(\alpha^{3/4}\epsilon)$, that is, $\hat{\eta}_c \rightarrow -\infty$ as $\alpha\epsilon^{4/3} \rightarrow 0$: the critical plane lies at $-\infty$ (we find $\hat{\eta}_v=-\infty$ as in the $\alpha \ll 1$ regime). However, once system (33)-(38) has been solved, equating the resulting heat flux at $\hat{\eta}=-\infty$, to the energy flux deposited at the critical plane,

$$-\frac{1}{16} \left[\frac{3}{4(Z_i+1)} \right]^{5/2} \hat{z}_e^{5/2} \frac{d\hat{z}_e}{d\hat{\eta}} \Big|_{\hat{\eta} \rightarrow -\infty} = \frac{Z_i+1}{12} \frac{\alpha^{1/2}}{\xi_f^{12/5}}$$

allows to determine ξ_f , and thus to complete the solution; we notice that system (33)-(38) leads indeed to $\hat{z}_e \sim (-\hat{\eta})^{2/7}$ as $\hat{\eta} \rightarrow -\infty$.

It may be shown analytically that in the neighborhood of $\hat{\eta}=0$, any solution to Eqs. (33)-(37) behaves as

$$\begin{aligned} \hat{z}_e &\approx A(-\hat{\eta})^{2/5}, & \hat{y} &\approx -\frac{A^{7/2}}{50B_1B_2} \left[\frac{3}{4(Z_i+1)} \right]^{5/2} (-\hat{\eta})^{2/5} \\ \hat{v} &\approx \frac{B_1B_2}{A} (-\hat{\eta})^{-2/5}, & \hat{z}_i &\approx \hat{z}_e - \left[\frac{3}{4(Z_i+1)} \right]^4 \frac{A^7 (-\hat{\eta})^{4/5}}{860b(B_1B_2)^2} \end{aligned}$$

where A is an arbitrary constant; it is possible to establish some bounds on A . Then, sweeping through the A -range one arrives at the value that allows satisfying condition (38). The resulting solution for the expansion flow is shown in Fig. 3 for $Z_i=1, 2$ and 4. We find that

$$\alpha^{-5/24} \xi_f = .39, .43, \text{ and } .49,$$

for $Z_i=1, 2$ and 4 respectively.

We finally notice that, as in I, there exists a very

thin transition layer centered at \bar{n} , where density peaks; its analysis would be similar to that carried out in the Appendix of I. We find $z=0(\alpha^{9/80})$, $v=0(\alpha^{-9/80})$ in this layer.

V. VALIDITY OF THE RESULTS

In the analysis, we assumed near-maxwellian distribution functions, that is, $\lambda_j \ll \Delta x$ and $t_j \ll t$, where λ_j and t_j are the mean free path and collision time for species j , Δx is the characteristic width of a particular region of the motion, and $t \ll \tau$. In self-similar variables the first condition reads

$$\theta_j^2 / \bar{n} \Delta \xi \ll (\alpha m_i / m_e)^{1/2}, \quad (39)$$

and is, naturally, equivalent to the requirement that no flux limiter be needed [see Eq. (9)]; the second condition, on the other hand, reads

$$\theta_e^{3/2} / \bar{n} \ll \alpha m_i / m_e, \quad (40)$$

$$\theta_i^{3/2} / \bar{n} \ll \alpha (m_i / m_e)^{1/2}, \quad (41)$$

for electrons and ions respectively.

For $1 \ll \alpha \ll \epsilon^{-4/3}$, the results of Sec. IV show that all three conditions (39)-(41) are well satisfied [implying in the more restrictive case, that $(m_i / m_e)^{1/2}$ be large]. For $\alpha \ll 1$, the stringest requirement is found to be (see Sec. III) inequality (39) for electrons, in both the thermal wave and expansion regions; for the thermal wave, (39) leads to

$$1 \ll (\alpha m_i / m_e)^{1/2},$$

while for the expansion we get

$$1 < \alpha (m_i/m_e)^{1/2},$$

an even stronger condition. Thus, the analysis of Sec. III for the expansion has a very limited range of validity, $(m_e/m_i)^{1/2} < \alpha < 1$. Clearly, for $\alpha \lesssim (m_e/m_i)^{1/2}$, the heat flux limiter given in Eq. (9) should be used. Actually, the flux limiter is even needed for the regimes $\alpha < 1$ and $1 < \alpha < \varepsilon^{-4/3}$, in the tail of the expansion where \bar{n} is small enough.

The hypothesis of quasineutrality, $\Delta x \ll \lambda_D$ (Debye length), used throughout the analysis, when written in self-similar variables reads

$$\omega \tau (t/\tau) \gg (A_i \varepsilon \alpha \theta_e / \bar{n})^{1/2} (10/\Delta \xi) ; \quad (42)$$

it is easily verified that for all reasonable values of ω , τ and A_i , (42) is well satisfied except at the very beginning of the pulse (or far in the expansion tail, where \bar{n} is extremely small).

VI. GENERAL DISCUSSION FOR ARBITRARY α

We are now in a position, using the results obtained here and I, to discuss how the plasma behaviour changes as α goes from large to small values. Figures 4a and b correspond to the regime $\alpha \gg \varepsilon^{-4/3}$ analyzed in I. In Fig. 4a we consider the case $\alpha \gg \varepsilon^{-5/3}$; there is then, beginning from the right, undisturbed plasma, a shock, a region of isentropic compression, a thin deflagration layer and finally a much wider isentropic expansion,

bounded by the vacuum at a finite distance from the origin ξ_v . The critical plane lies in the deflagration layer, to the right of the origin ($\xi_c > 0$). In Fig. 4b we consider the case $\varepsilon^{-4/3} \ll \alpha \ll \varepsilon^{-5/3}$, the only changes with respect to Fig. 4a, being that now the deflagration is much wider than the compression region and ξ_c is negative. Clearly, if $\alpha \sim \varepsilon^{-5/3}$, the compression flow and the deflagration layer have the same size.

Figure 4c for $\alpha \sim \varepsilon^{-4/3}$ may be inferred as intermediate limit from the results obtained in I and in Sec. IV. The main change is the merging of the deflagration and expansion regions (which thus ceases to be isentropic); the critical plane lies at a finite distance within it. Figure 4d corresponds to the regime $\varepsilon^{-4/3} \gg \alpha \gg 1$ and shows schematically results from Sec. IV, the only change being that the plasma vacuum boundary lies at infinity (the density within the expansion region is much larger than the critical density). We notice that throughout Figs. 4a-d there is a transition layer (where temperature is very low) just to the left of the isentropic compression, and an electron precursor ahead of the shock,⁷ both being very thin.

Figure 4e for $\alpha \sim 1$ may be inferred as intermediate limit from the results of Secs. III and IV. As α decreases to values of order unity the precursor thickness and the temperature minimum grow until the temperature is of the same order everywhere and the shock stands in the middle of a quasi-thermal wave (convection being important, although there is isentropic flow nowhere).

As α goes on decreasing, the shock moves to the origin and its intensity weakens. Figure 4f shows finally the regime $\alpha \ll 1$; the shock has become a weak discontinuity close to the origin, convection being negligible ahead of it. The expansion flow has become isothermal (electron conduction being dominant) and thin compared with the thermal wave. Throughout Figs. 4d-f the critical plane lies far in the expansion tail which reaches infinity. Actually as noticed in Sec. V, the plasma is collisionless in the far tail, and therefore the validity of the results breaks down there.

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